A R T I C L E

CHARMED BARYONS

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Abstract

A review of the spectroscopy and decay properties of charmed baryons is presented, with emphasis on the double–charm and triple–charm sectors, which have been studied recently.

1. Introduction

The discovery of hidden charm in 1974 and naked charm in 1976 opened a new era in hadron spectroscopy [1, 2]. All new heavy mesons and baryons [3] have been described simply by adding a new quark, c, to the series of the constituents d, u and s, thus reinforcing the quark model.

This qualitative success stimulated detailed dynamical calculations of quark binding, based on models of confinement. Several bag or potential models were very efficient in reproducing the static properties of hadrons. Meanwhile, some progress has been made in non–perturbative QCD, thanks, for instance, to lattice simulations, which provide some basis for empirical models. In this framework, we shall show that baryons with single, double, or triple charm provide very interesting opportunities to test the dynamics of quark confinement.

In charmonium, the heavy quarks move slowly in the short–range QCD potential, and form several narrow levels, with interesting cascades from one to another. In charmed mesons, the light antiquark is accelerated by the presence of a heavy static colour source, so that relativistic effects are enhanced.

These two aspects of quark dynamics also enter in the charmed baryon sector, where they are sometimes intimately combined. Triple–charm baryons are the analogues of charmonia, with even narrower levels and the subtleties of the three–body systems. In double–charm baryons, we see the slow relative motion of the two heavy quarks together with the relativistic motion of a light quark around a coloured centre. Single–charm baryons look more like charmed mesons, but some details such as hyperfine structure allow one to disentangle the heavy–light from the light–light type of interquark forces.

It is hardly necessary to underline the importance of charm for weak interactions. The charmed quark was in fact predicted by GIM [4] to solve a long–standing problem related to the decay properties of neutral kaons. Charmed particles themselves have revealed intriguing decay properties: in particular, the ratio of meson lifetimes, $\tau(D^+)/\tau(D^0)$, is different from 1, showing that the decaying charmed quark does not ignore its environment. The dynamics of the weak decay of charmed mesons nowadays seems better understood [5]. We shall show that charmed baryons permit us better to determine the relative importance of W emission, W exchange and final–state interactions in weak decays [6, 7].

Throughout this review, we shall very often use the convenient language of the non–relativistic quark model. Most of the physics discussion on charmed baryons holds, however, beyond this particular approximation. We shall give references to alternative studies based on bag models or relativistic wave equations, and emphasize the need for relativistic corrections when discussing electromagnetic mass differences.

2. The experimental situation

We adopt here the notation of the Particle Data Group [1] which now replace the naming scheme of ref. [8]: Λ_c, for instance is a Λ (sud) state where the strange quark s has been replaced by a c quark. In the (csq) sector ($q = u$ or d), we denote by Ξ_c the baryon where the (sq) pair is most often in a spin 0 state, Ξ'_c the baryon with total spin 1/2 and the (sq) pair mostly in a spin 1 state, Ξ^*_c the lowest spin 3/2 baryon, and Ξ^{**}_c any higher radial or orbital excitation. Note that Ω_c is a spin 1/2 baryon with (css) content, while Ω^*_c is its spin 3/2 partner. The ground states of charmed mesons and baryons are listed in table 1.

Table 1 Ground–state hadrons with charm (q denotes u or d quark)

	Charmed mesons				Charmed baryons												
Quark content	$c\bar{q}$		$c\bar{s}$		cqq	cqq		csq			ccq		css		ccs		ccc
Isospin	1/2		0		0	1		1/2			1/2		0		0		0
Spin	0	1	0	1	1/2	1/2	3/2	1/2	1/2	3/2	1/2	3/2	1/2	3/2	1/2	3/2	3/2
Name	D	D^*	D	D^*_s	Λ_c	Σ_c	Σ^*_c	Ξ_c	Ξ'_c	Ξ^*_c	Ξ_{cc}	Ξ^*_{cc}	Ω_c	Ω^*_c	Ω_{cc}	Ω^*_{cc}	Ω_{ccc}

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The experimental masses and lifetimes of charmed baryons are summarized in table 2. These baryons have been seen in several experiments, using hadronic or leptonic beams. There is an urgent need for high–precision and high–statistics measure-ments of these baryons. This will hopefully be achieved in future experiments with hyperon beams or at flavour factories.

Table 2
Experimental masses and liftetimes of charmed baryons

Baryon	Mass (GeV)	$\tau\,(10^{-13}\,\mathrm{s})$	Ref.
Λ_c	2.285	(1.8 ± 0.2)	[1]
Σ_c	2.455	unstable ($\to \Lambda_c + \pi$)	[1]
Ξ_c^+	2.460	(4.3 ± 1.5)	[1]
Ξ_c^0	2.471		[2]
Ω_c^0	2.740		[1]

3. Charmed baryon spectroscopy

3.1 Baryon mass inequalities

In QCD, gluons are coupled to the colour rather than to the flavour of quarks. The interquark potential is thus, in first approximation, flavour–independent. Likewise, in QED, the same electrostatic potential e^2/r binds e^+e^- as well as e^-p. Flavour independence leads to interesting inequalities among hadrons.

Let $\mathcal{M}(m_1, m_2)$ be the mass of a $(q_1\bar{q}_2)$ meson in its ground state. Neglecting spin corrections, one gets [9]

$$\mathcal{M}(M, M) + \mathcal{M}(m, m) \leq 2\,\mathcal{M}(M, m) \qquad (1)$$

from the variational principle or the observation that the reduced Hamiltonian depends linearly on the inverse reduced mass λ and that, if $H = A + \lambda B$, its ground–state energy is a concave function of λ. The inequality (1) can be extended to the sum of the first n levels or applied as it stands to the lowest state in a sector of a given parity or orbital angular momentum. Applying the inequality (1) to charmed and strange quarks gives

$$\mathcal{M}(J/\psi) + \mathcal{M}(\phi) \leq 2\mathcal{M}(Ds) \qquad (1')$$

which is indeed true [1].

In a first generalization to baryons, one may study the ground–state binding energy $E_0(m_1, m_2, m_3)$ as a function of m^{-1}, the inverse mass of the first quark, and constrain the mass of Λ_b in terms of the masses of Λ_c and Λ [10]. However, the resulting inequalities for baryons depend on the quark masses.

On the other hand, there are simple convexity relations involving the three inverse masses and leading to inequalities such as

$$\mathcal{M}(m_1,m_1,m_1) + \mathcal{M}(m_2,m_2,m_2) + \mathcal{M}(m_3,m_3,m_3) \leq 3\,\mathcal{M}(m_1,m_2,m_3)\ . \qquad (2)$$

This provides a limit on the mass of (ccc) from that of (qqq), (sss) and the recently discovered (csq) [1, 2]. With a plausible estimate of hyperfine corrections for (csq), one obtains

$$\mathcal{M}(ccc) \leq 5.08 \text{ GeV} \ . \qquad (3)$$

The inequality (1) holds for any interquark potential, provided it does not depend on flavour. The following inequality like (2), requires that the potential does not grow too sharply [11], a condition easily fulfilled by realistic models. It reads

$$\mathcal{M}(M, M, m) + \mathcal{M}(m, m, m) \leq 2\mathcal{M}(M, m, m) \ , \qquad (4)$$

and gives a limit on double–charm baryon masses. When hyperfine corrections are omitted, one obtains

$$\mathcal{M}(ccq) \leq 3.72 \text{ GeV} \ . \qquad (5)$$

3.2 Inequalities relating mesons and baryons

There are good reasons to believe that the potential that links the three quarks of a baryon can be simulated by a sum of two–body terms, each of these being half of the quark–antiquark potential binding mesons [12]

$$V(r_1, r_2, r_3) = \sum v_{qq}(r_{ij}) = \frac{1}{2}\sum v_{q\bar{q}}(r_{ij}) \ . \qquad (6)$$

If one takes seriously this "1/2 rule", one gets for equal–mass quarks [13]

$$2\,\mathcal{M}(QQQ) > 3\mathcal{M}(Q\bar{Q}) \qquad (7)$$

as a direct consequence of the variational principle. This means that a quark feels itself heavier in a baryon than in a meson(*). If the potential of eq. (6) is considered as acting between quarks with parallel spins, one may apply the inequality (7) between a spin 3/2 baryon and a vector meson. In the strangeness sector, one checks that $2\,\mathcal{M}(\Omega^-) > 3\,\mathcal{M}(\phi)$. In the charm case, one predicts

$$\mathcal{M}(\Omega_{ccc}) > \frac{3}{2}\mathcal{M}(J/\psi) \approx 4.65 \text{ GeV} \ . \qquad (8)$$

Various generalizations of the inequality (7) can be elaborated, involving the sum of the first levels, or different flavours, or spin 1/2 baryons. If one neglects spin effects or considers only parallel spins, one may write [14]

$$2\,\mathcal{M}(Q_1Q_2Q_3) > \mathcal{M}(Q_1\bar{Q}_2) + \mathcal{M}(Q_2\bar{Q}_3) + \mathcal{M}(Q_3\bar{Q}_1)\ , \qquad (9)$$

whereas incorporating hyperfine effects for spin 1/2 baryons and for mesons, in the single–charm sector, leads to [10]

(*) In the Coulomb case, in a very different context, an equivalent equality is discussed by J.M. Lévy–Leblond (ref. [13]). We thank A. Martin for an interesting discussion on this point.

$$\mathcal{M}(\Lambda_c) > \frac{1}{2}\mathcal{M}(\pi) + \frac{3}{4}\mathcal{M}(D^*) + \frac{1}{4}\mathcal{M}(D)\,(2.28 > 2.04\ \text{GeV})$$
$$\mathcal{M}(\Sigma_c) > \frac{1}{2}\mathcal{M}(\rho) + \frac{3}{4}\mathcal{M}(D) + \frac{1}{4}\mathcal{M}(D^*)\,(2.45 > 2.29\ \text{GeV}) \quad (10)$$
$$\mathcal{M}(\Sigma_c^*) > \frac{1}{2}\mathcal{M}(\rho) + \mathcal{M}(D^*) \approx 2.40\ \text{GeV} \ .$$

For double–charm baryons, one predicts

$$\mathcal{M}(\Xi_{cc}) > \frac{1}{2}\mathcal{M}(J/\psi) + \frac{3}{4}\mathcal{M}(D) + \frac{1}{4}\mathcal{M}(D^*) \approx 3.45\ \text{GeV} \quad (11)$$
$$\mathcal{M}(\Xi_{cc}^*) > \frac{1}{2}\mathcal{M}(J/\psi) + \mathcal{M}(D^*) \approx 3.56\ \text{GeV} \ .$$

The case of baryons with charm and strangeness is discussed in ref. [10].

One should mention that some generalizations of the inequality (7) do not always hold. One gets, for instance

$$\mathcal{M}(QQQ) + \mathcal{M}(\bar{q}\bar{q}\bar{q}) < 3\,\mathcal{M}(Q\bar{q}) \quad (12)$$

if the mass ratio M/m is large enough. This means that the antibaryon $\bar{\Omega}_{ccc}$ made out of three charmed antiquarks does not annihilate when it touches ordinary matter.

3.3 Single–charm baryons

One can hardly test in detail the central potential in the (Qqq) sector without data on orbital and radial excitations. One can check, however the validity of flavour independence. Using a reasonable value for the mass of the charm quark, one can reproduce the masses of Λ_c, Σ_c and Ξ_c with a potential which fits other ground states such as N, Δ, Λ, or Ω^- [3, 15].

In fact, most studies on charmed baryons were devoted to hyperfine splittings, to test current ideas on spin–dependent forces and to predict whether spin excitations such as Σ_c, Σ_c^*, Ξ'_c or Ξ_c^* are stable or not.

A success of the quark model is the systematic description of hyperfine splittings with a chromomagnetic interaction [16]

$$V_{ss} = \frac{C}{2}\sum_{i<j}\frac{\sigma_i\sigma_j}{m_i m_j}\,\delta^{(3)}(r_{ij}) \ . \quad (13)$$

The $(m_i m_j)^{-1}$ is crucial, especially in explaining the Σ–Λ mass difference. In the Λ, the chromomagnetic attraction occurs entirely between light quarks, and is proportional to m_q^{-2} whereas in Σ, there are terms in $(m_q m_s)^{-1}$ involving strange and ordinary quarks. If the s quark is now replaced by c or another heavy quark Q, one obtains [17]:

(a) $(\Sigma_Q–\Lambda_Q)$ grows with m_Q and quickly becomes larger than the crucial threshold of 140 MeV ($= m_\pi$). As $m_Q \to \infty$, this quantity becomes constant. Experimentally, $\Sigma_c - \Lambda_c \approx 170$ MeV [1].

(b) On the other hand, $(\Sigma_Q^* - \Sigma_Q)$ vanishes as $m_Q \to \infty$. Indeed, in the model (13), this splitting results only from the interaction between a light and a heavy quark.

The analysis is slightly more involved in the sector with charm and strangeness and the predictions are slightly more model-dependent. The $\Xi'_c - \Xi_c$ splitting is always found to be less than 140 MeV, so that the Ξ'_c should be rather narrow. On the other hand, the transition $\Xi_c^* \to \Xi_c + \pi$ is either easily [18] or marginally [19] allowed, so that the Ξ_c^* is either broad or relatively narrow, depending on the model. The spin excitation Ω_c^* should decay radiatively into the ground state Ω_c.

One model for the description of charmed baryons was proposed in ref. [18] based on the central potential

$$V = \frac{1}{2}\sum_{i<j}\left(A + B r_{ij}^{\beta}\right) , \quad (14)$$

supplemented by a spin–spin term (13), treated to first order. With the following parameters, given in GeV or appropriate powers of GeV, $m_q = 0.300$, $m_s = 0.600$, $m_c = 1.895$, $A = -8.337$, $B = 6.9923$, $\beta = 0.1$, and $C = 2.572$, one gets the results shown in table 3.

Table 3
Ground–state baryons from the potential model of eqs (13) and (14). Masses are in GeV

Baryon	N	Δ	Λ	Σ	Σ^*	Ξ	Ξ^*	Ω^-
Theory	input	input	1.111	1.176	1.392	1.304	1.538	input
Exp.	0.938	1.232	1.115	1.193	1.383	1.318	1.533	1.672
Baryon	Λ_c	Σ_c	Σ_c^*	Ξ_c	Ξ'_c	Ξ_c^*	Ω_c	Ω_c^*
Theory	input	2.443	2.542	2.457	2.558	2.663	2.664	2.775
Exp.	2.282	2.455		2.460				

The fit of the horizontal spectrum is obviously satisfactory. It is worth stressing the ideological evolution through the years. The quark model came out of the flavour symmetry $SU(3)_F$, which combines isospin and strangeness. The generalization to $SU(4)_F$, to include charm, is straightforward but is not very useful, since the breaking of symmetry is too great with, for instance, the Δ (1.2 GeV) and Ω_{ccc} ($\cong$ 4.5 GeV) lying in the same multiplet. The true symmetry is associated with the colour group $SU(3)_c$, which leads to the flavour independence of the confining potential and to the simple chromomagnetic corrections (13).

3.4 Double–charm baryons

A detailed study of the spectroscopy of baryons with double charm has recently been carried out [7]. The internal dynamics of these hadrons is rather interesting. While the two heavy quarks move slowly and remain close to each other, the light quark oscillates at quite a speed far away from the charmed centre. To some extent, the double–charm baryons combine the internal structures of charmonium and charmed mesons.

In their ground state, double–charm baryons exhibit a well–pronounced quark–diquark (q–cc) structure, since the average distance between the two heavy quarks is much smaller than the distance between them and the light quark. However, a quark–diquark approximation with a frozen diquark does not account for the excitation spectrum. It is more economical, indeed, to promote the radial or orbital motion of the heavy quarks than the motion of the light quark. The first levels of double–charm baryons correspond essentially to internal excitations of the diquark.

In fact, an efficient approximation is provided by the Born–Oppenheimer method. For a given separation of the charmed quarks, one computes the binding energy of the light quark which feels the (non–central) potential of the two coloured sources. This binding energy, when supplemented by the direct c–c interaction, provides an effective potential which governs the spectrum and the wave–function of the two–charmed quarks. As shown in ref. [7] and in many other examples in atomic or nuclear physics, the Born–Oppenheimer approximation works astonishingly well when compared to the exact solution of the three–body problem.

In ref. [7], the Born-Oppenheimer method was tested in a non–relativistic model, for which an exact three–body calculation is feasible using, for instance, the method of hyperspherical co–ordinates. The Born–Oppenheimer method can also be used with the bag model [20], where the light quark is treated relativistically [7]. Here, as in the charmonium case [21], the bag is used not for calculating directly the hadron mass, but for providing the effective potential between the heavy quarks.

The masses of double–charm baryons have been computed in ref. [7] using various potentials and bag models. In the latter case, the results are rather sensitive to details of the calculation such as centre–of–mass correction, zero–point energy, or running coupling constant. In fact, the simpler potential models give more convergent results, suggesting a better extrapolating power. To guide the discussion, we display in table 4 the results obtained with the potential (13)–(14).

Table 4

Masses of double–charm baryons, obtained from the potential model (13)–(14). For radial (n = 1) or orbital excitations, hyperfine corrections are neglected. Units are in GeV

	$n = \ell = 0$		$n = 1$	$n = 0$
	$s = 1/2$	$s = 3/2$	$\ell = 0$	$\ell = 1$
ccq	3.613	3.741	4.110	3.971
ccs	3.703	3.835	4.235	4.084

Some comments are in order:

(a) The spin excitation Ξ_{cc}^{*} cannot decay into $\Xi_{cc} + \pi$ and hence should be rather narrow.

(b) On the other hand, the radial or orbital excitations of (ccq) are very unstable, thanks to the decay $\Xi^{**} \to \Xi_{cc} + \pi$.

(c) The spin excitation Ω_{cc}^{*} and the first orbital excitation Ω_{cc}^{**} of the (ccs) system are stable under the strong interactions and should decay radiatively.

(d) Starting with the first radial excitation of (ccs), reactions like $\Omega_{cc}^{**} \to \Xi_{cc} + K$ can occur, resulting in large widths.

3.5 Baryons with triple–charm

As underlined by Bjorken and Martin [22], the (ccc) spectrum is one of the most exciting goals in baryon spectroscopy. We are dealing here with a simple system, almost non–relativistic, with several narrow or rather narrow levels. The spectrum will provide direction information on the QCD potential between three coloured charges. The (ccc) system deserves a detailed treatment of the three–body problem, beyond the approximation of a perturbed oscillator. Methods are available [16].

One of the interesting problems that (ccc) spectroscopy will clarify, is the following. With any plausible local potential, the first orbital excitation, with negative parity, always lies below the first radial excitation which has the same quantum numbers as the ground state [23]. This generalizes a well–known theorem in the two–body case [24]. For ordinary baryons, the experimental candidate for radial and orbital excitations are almost degenerate, leading to many discussions of relativistic effects, gluon correc-tions, coupling to decay channels, etc. It seems important to check that for the (ccc) case, one gets a normal ordering, other-wise our understanding of baryons would have to be revised. Other problems concern the breaking of the N = 2 and N = 3 levels of the harmonic oscillator [25], spin–orbit forces [26], etc.

Let us discuss briefly the stability of the (ccc) levels. The first excitations lie below the threshold for $\Omega_{ccc}^{**} \to (ccq) + (\bar{q}c)$ which plays the same role as the $D\bar{D}$ threshold for charmonium. The first (ℓ = 1) states should decay radiatively into the ground state at a slow rate, however, since the Pauli principle requires that the initial state has spin 1/2 and final state spin 3/2. Also of interest are the isospin–violating reactions $(ccc)^{***} \to (ccc) + \pi^0$, to be compared to $\psi' \to J/\psi + \pi^0$. Some transitions induced by the emission of an η meson should also be seen.

To end this section, we compare in table 5 three estimates of the first levels of (ccc). The first one results from the "1/2 rule" of eq. (6) applied to a popular quarkonium potential [12]. For the two others, $\mathcal{M}(QQQ)$ has been calculated using the adiabatic approximation to the bag model [21]. The results are quite similar, and compatible with the inequalities (3) and (8).

Table 5

Predictions for the masses of the first (ccc) levels, in GeV

	Richard [12]	Hasenfratz et al. [21]	Aerts et al. [21]
$L = 0, n = 0$	4.797	4.791	5.040
$L = 1, n = 0$	5.112	5.137	5.305
$L = 0, n = 1$	5.278	5.302	5.430

4. Breaking of isospin symmetry

The mass shifts between hadrons of the same isospin multiplet have often stimulated interesting discussions. Much effort has been devoted to understand the neutron–proton mass difference m_n–m_p. In the constituent quark model, it is due to the electromagnetic interactions between quarks, to the mass difference between the u and d quarks, and to the resulting difference of their binding energy and chromomagnetic interaction [27].

These various effects sometimes tend to cancel, so that the net effect is rather small and sensitive to the details of the calculation of each term. For instance, in the charm sector, Lane and Weinberg [28] have predicted much smaller splittings than De Rújula, Georgi and Glashow [29] who use a simple non–relativistic model to extrapolate matrix elements from ordinary to charmed hadrons. The recent measurement by the CLEO Collaboration [30] shows that, indeed Σ_c^0 and Σ_c^{++} are almost degenerate.

One recalls here that the non–relativistic approximation is never worse than for open flavours (likewise, the electron is more relativistic in hydrogen than in protonium). If one computes the mass difference $\mathcal{M}(D^+) - \mathcal{M}(D^0)$ or the analogue for charmed baryons by accounting only for the quark–mass difference $\delta_m = m_d - m_u$ and neglecting provisionally any electromagnetic or hyperfine effects, one finds that the change in binding energy is larger, in absolute value, than δ_m [31]. In other words, the hadron mass decreases when the mass of one of the constituents is increased. This suggests, perhaps better than any estimate of average velocities, the need for a relativistic treatment.

5. Weak decays of charmed baryons

When it was announced that the ratio $R = \tau(D^+)/\tau(D^0)$ of the lifetimes of charged and neutral charmed mesons is substantially larger than 1, this was a great surprise, especially because the first measurements had exaggerated the effect. The present value [1]

$$R = \frac{(10.7 \pm 0.3)\times 10^{-13}\text{s}}{(4.3 \pm 0.1)\times 10^{-13}\text{s}} \approx 2.5 \qquad (15)$$

remains quite impressive. Also of interest is the observation that $\tau(D_s) \approx \tau(D^0)$ (4.4 ± 0.3 vs 4.3 ± 0.1 in units of 10^{-13} s). Everyone is now convinced that the charmed–quark decay depends on the environment. However, the relative importance of the various mechanisms remains an open question.

The leading contributory mechanism is certainly the spectator diagram of fig. 1. In a first possible scenario, this spectator diagram dominates over other weak mechanisms and the rate for D^+ decay is decreased by the Pauli principle or, say, by the interference between the spectator $\bar{d}$ and the $\bar{d}$ produced by the hadronic mode $W^+ \to u\bar{d}$. This model explains why $\tau(D_s) \approx \tau(D^0)$, but provides a ratio R that is too small. It also leads to a semileptonic branching ratio for D^+ that is too large, whereas the experimental value for $BR\ (D^+ \to e^+ \nu X)$ is (19 ± 2)%, in agreement with a naive fermion counting in fig. 1.

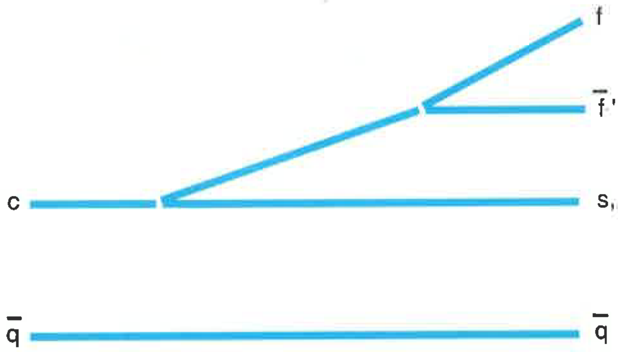

FIGURE 1

Spectator diagram for the decay of charmed mesons.

Other explanations rely more on the variety of weak mechanisms than on final–state interactions. The D^0 benefits from the exchange diagram of fig. 2, leading to a rather short lifetime. This diagram does not apply to D^+ or D_s.

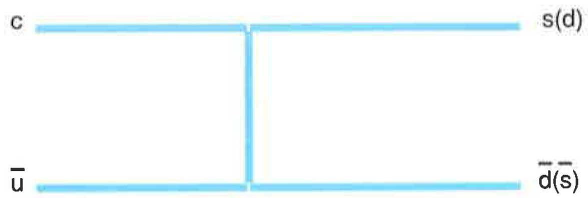

FIGURE 2

W–exchange diagram.

The $D_s = (c\bar{s})$ lifetime is shortened by the annihilation diagram of fig. 3, which does not contribute much for D^+, since the $c\bar{d} \to W^+$ coupling is suppressed in the Cabibbo–Kobayashi–Maskawa matrix.

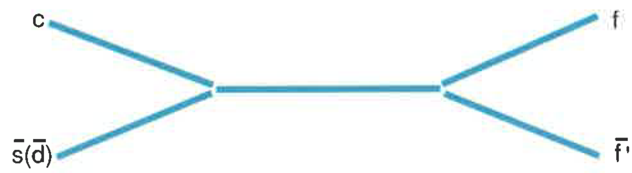

FIGURE 3

Annihilation diagram for D_s or D^+ decay.

Let us finally mention that the penguin diagrams (fig. 4) are not very important in the charm case, and, anyhow, do not directly influence the ratio $R = \tau(D^+)/\tau(D^0)$.

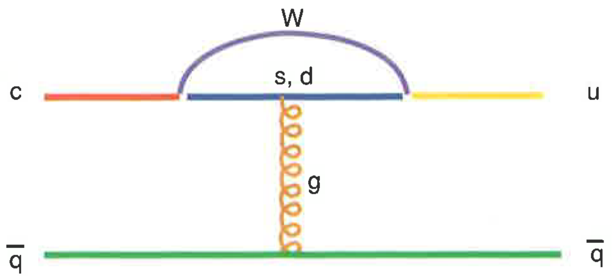

FIGURE 4

Penguin diagram.

As underlined by several authors [5], more data are needed, especially on D_s and some rare decay modes, to measure in detail the role of each mechanism. Charmed baryons will certainly provide very valuable information since the weight of some contributions is modified with respect to the meson case. For instance, the exchange diagram of fig. 2 does not suffer from helicity suppression. On the other hand, the annihilation diagram of fig. 3 does not contribute, if one restricts the calculation to valence quarks.

The case of single–charm baryons was analysed by Guberina et al. [6], who ended with the hierarchy of lifetimes

$$\tau(\Omega_c^0) \leq \tau(\Xi_c^0) < \tau(\Lambda_c) < \tau(\Xi_c^+) \ . \qquad (16)$$

The Pauli principle plays the leading role here, with a constructive interference between s quarks for Ξ_c^0 et Ω_c^0. Next comes the exchange diagram, which favours Λ_c with respect to Ξ_c^+.

When a similar analysis is carried out for baryons with double charm, one gets the ordering [7]

$$\tau(\Xi_{cc}^+) < \tau(\Omega_{cc}^+) < \tau(\Xi_{cc}^{++}) \ , \qquad (17)$$

which deserves experimental testing. The Ξ_{cc}^{++} decay is penalized by the interference between the spectator u quark and the u quark due to the hadronic mode of W. The Ω_{cc}^+ decay benefits from the constructive interference between the spectator s quark and the s quark resulting from the decay of one of the charmed quarks. Finally, Ξ_{cc}^+ is helped by the W–exchange diagram.

6. Conclusion

Charmed baryons reveal fundamental aspects of quark confinement and weak decay and, without doubt, future experiments will provide us with more information on this physics.

In the past, hyperon physics turned out to be very instructive. In spectroscopy, many debates, with key questions, resulted from the measurements of the Σ–Λ or $\Lambda(1520)$–$\Lambda(1405)$ mass differences. Charmed baryons will certainly give us some surprises.

Also, strangeness deposition in nuclei stimulated many studies on hypernuclei, where the baryon–baryon interaction, the baryon–nucleus interaction and the possible renormalization of strong forces and weak decays in the nuclear medium all play a role. Charmed nuclei have already been discussed [32]. The charmed baryons could go deeper inside the nucleus than hyperons and, perhaps, feel some partial deconfinement.

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References

[1] Review of Particle Properties, Phys. Lett. B 204 (1988) 1, (for experimental results, and references to first measurements).
[2] P. Avery et al., Phys. Rev. Lett. 62 (1989) 863.
[3] W. Kwong, J.L. Rosner and C. Quigg, Annu. Rev. Nucl. Part. Sci. 37 (1987) 325 (for a review on charm);
T. Appelquist et al., Annu. Rev. Nucl. Part. Sci. 28 (1978) 387.
[4] S.L. Glashow, J. Iliopoulos and L. Maiami, Phys. Rev. D7 (1970) 1285.
[5] I.I. Bigi, Proc. of the 14th Annual SLAC Summer Institute on Particle Physics, Stanford (1986);
M. Wirbel, Progress in Nucl. and Part. Physics, ed. A. Faessler, 21 (1988) 33–98;
B. Eisenstein, Physics in Collision 6, ed. M. Derrick, World Scientific 1986;
J.L. Basdevant, I. Bedagia and E. Predazzi, Nucl. Phys. B294, (1987) 1054.
[6] B. Guberina, R. Rückl and J. Trampetic, Zeitschr. für Phys. C33 (1986) 297.
[7] S. Fleck, Thèse de l'Université Joseph Fourier, Grenoble (1988).
S. Fleck and J.M. Richard, Progr. in Theor. Phys. 82 (1989) 760.
[8] M.K. Gaillard, B.W. Lee and J.L. Rosner, Rev. Mod. Phys. 47 (1975) 277.
[9] R.A. Bertlmann and A. Martin, Nucl. Phys. B168 (1980) 111;
S. Nussinov, Phys. Rev. Lett. 52 (1984) 966.
[10] A. Martin, Phys. Lett. B103 (1981) 51;
A. Martin and J.M. Richard, Phys. Lett. B185 (1987) 426.

References (Cont'd)

[11] E.H. Lieb, Phys. Rev. Lett. 54 (1985) 1987;
A. Martin J.M. Richard and P. Taxil, Phys. Lett. 176B (1986) 224.

[12] O.W. Greenberg and H.J. Lipkin, Nucl. Phys. A370 (1981) 349;
D.P. Stanley and D. Robson, Phys. Rev. Lett. 45 (1980) 235;
J.M. Richard, Phys. Lett. 100B (1981) 515.

[13] J.P. Ader, J.M. Richard and P. Taxil, Phys. Rev. D25 (1982) 2370;
J.M. Lévy–Leblond, J. Math. Phys. 10 (1969) 806.

[14] S. Nussinov, Phys. Rev. Lett. 51 (1983) 2081;
J.M. Richard, Phys. Lett. B139 (1984) 408.

[15] J.L. Basdevant and S. Boukraa, Zeitschr. für Phys. C28 (1985) 413;
S. Capstick and N. Isgur, Phys. Rev. D34 (1986) 2809.

[16] A. De Rújula, H. Georgi and S.L. Glashow, Phys. Rev. D12, (1975) 147;
H.J. Lipkin, Baryon '80, Proc. of the IVth Int. Conf. on Baryon Resonances, Toronto, ed. N. Isgur (1980).

[17] J.M. Richard and P. Taxil, Ann. Phys. (NY) 150 (1983) 267.

[18] J.M. Richard and P. Taxil, Phys. Lett. 128B (1983) 453.

[19] K. Maltman and N. Isgur, Phys. Rev. D22 (1980) 1701.

[20] C.E. De Tar and J.R. Donoghue, Annu. Rev. Nucl. Part. Sci. 33 (1983) 235.

[21] P. Hasenfratz and J. Kuti, Phys. Rep. 42 (1978) 75;
L. Heller, in Quarks and Nuclear Forces, ed. D.C. Fries and B. Zeitnitz, Springer–Verlag (1982) 145;
P. Hasenfratz et al., Phys. Lett. 94B (1980) 401 and 95B (1980) 299;
A.T. Aerts and L. Heller, Phys. Rev. D25 (1981) 1365.

[22] B.J. Bjorken, Proc. of the Int. Conf. on Hadron Spectroscopy, College Park, ed. S. Oneda, AIP (1985);
A. Martin, Proc. of the Workshop on Heavy Flavours, Erice 1988, ed. L. Cifarelli, Plenum, NY (1989).

[23] H. Høgaasen and J.M. Richard, Phys. Lett. 124B (1985) 520.

[24] H. Grosse and A. Martin, Phys. Rep. 60 (1980) 341.

[25] D. Gromes and I.O. Stamatescu, Nucl. Phys. B112 (1976) 213 and Zeitschr. für Phys. C3 (1979) 43;
N. Isgur and G. Karl, Phys. Rev. D19 (1979) 2653 and D23 (1981) 817;
R.H. Dalitz et al., J. Phys. G 3 (1977) L195;
K.L. Bowler et al., Phys. Rev. D24 (1981) 197;
J.M. Richard and P. Taxil, CERN–TH 5272–89 (to be published in Nucl. Phys. B).

[26] L.J. Reinders, Baryon '80, Proc. of the IVth Int. Conf. on Baryon Resonances, Toronto, Ed. N. Isgur (1980).

[27] N. Isgur, Phys. Rev. D21 (1980) 779.

[28] K. Lane and S. Weinberg, Phys. Rev. Lett. 37 (1976) 717.

[29] A. De Rújula, H. Georgi and S.L. Glashow, Phys. Rev. Lett. 37 (1976) 398.

[30] T. Bowcock et al., Phys. Rev. Lett. 62 (1989) 1240.

[31] W. Celmaster, Phys. Rev. Lett. 34 (1976) 1042;
J.M. Richard and P. Taxil, Zeitschr. für Phys. C26 (1984) 421.

[32] C.B. Dover and S.H. Kahana, Phys. Rev. Lett. 39 (1977) 1506.
R. Gatto and F. Paccanoni, Nuovo Cimento 46A (1978) 313.

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